LOWER HYBRID SOLITON MEDIATED RADIO EMISSION

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Abstract

A direct plasma emission process of radio waves is proposed. It is based on the assumption of acceleration of electrons in the combined electrostatic fields of high frequency (upper hybrid) waves and low frequency (lower hybrid) envelope solitons. Various frequency bands are excited in this process by the mixing of the wave components. The most intense power is obtained for emission in the X–mode near the fundamental frequency $\omega_h + \omega_l$ where h, l correspond to high and low frequencies, respectively. The emitted energy is taken from the high frequency field, while it is adiabatically modulated by the low frequency field. It is collected from a large volume of high frequency waves crossed by the soliton in a kind of vacuum cleaner effect. Since emission is restricted to the volume of the soliton, the efficiency is higher than in wave coupling, but to explain AKR intensity, about 10^6 solitons are required in the source region corresponding to a filling factor of about 1–10%. For slightly oblique propagation of the lower and upper hybrid waves the frequency of the wave is above the X–mode cut–off.

1 Introduction

Nonthermal radio emission from natural plasmas as observed to originate from planetary magnetospheres, the Sun and other weakly magnetized stars, can be generated by a number of processes which only require the presence of free energy. The generation of radio waves is obviously one very efficient way to get rid of the excess energy and transmit it into free space. The free energy is usually stored in forms like an active magnetic field configuration which can undergo reconnection and generate nonthermal particle distributions: Shock waves, electron beams, loss cone distributions etc.

Two different classes of nonthermal radiation processes are known. The most important one is the maser mechanism. It is basically linear and hence very efficient. It requires

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the presence of at least slightly relativistic electrons and of an anisotropic velocity space distribution with the excess energy in the perpendicular component. Such distributions are easily generated in a loss–cone magnetic mirror field configuration. As has first been recognized by Wu and Lee [1979] such a distribution may drive directly unstable electromagnetic waves propagating in the X–mode thereby decreasing the anisotropy. Since the waves in the source region have low group velocities, there is time enough to reach quasilinear equilibrium and to partially stabilize the distribution function. Conditions of this kind are realized in the auroral magnetosphere where they may cause the observed [Gurnett, 1974; Benson and Calvert, 1979] auroral kilometric radiation (AKR) as well as in other planetary magnetospheres.

In other cases when there is no loss cone available as for instance in the case of solar type III bursts, other and nonlinear mechanisms become important. Some of these mechanisms are known under the term of nonlinear wave coupling processes [Melrose, 1991]. They require the resonant collision of three waves, in general two electrostatic plasma waves and one electromagnetic wave which is the escaping radio wave. The efficiency of these processes is in general proportional to the product of the energy densities of the two participating electrostatic wave modes, and because of the resonance conditions (conservation laws) to be satisfied, the wave numbers of the electrostatic modes have to approximately cancel to yield the long wave length of the electromagnetic mode while the energy of the escaping mode becomes the approximate sum of the two participating electrostatic wave energies. This means that the escaping wave frequency is the sum of the two electrostatic wave frequencies. To determine the wave amplitudes in this process one has to solve the wave kinetic equation of the interaction. It is clear that this higher order process is not very efficient. Nevertheless it is believed to be of some importance in type III and type II burst radiation of the Sun as well as in type III burst emissions in front of the Earth's bow shock wave [Treumann et al., 1986; Lacombe et al., 1985]. Surprisingly enough it has never been observed in the auroral magnetosphere where precipitating auroral electron beams are frequently present.

2 The model

When the electrostatic waves involved possess large growth rates they can readily enter into the nonlinear regime and under certain circumstances can evolve into solitons or cavitons. The electrostatic waves become localized, and the interaction is restricted to a small volume in space. This may in some cases be favorable for the emission process [Papadopoulos and Freund, 1978; Treumann and Bernold, 1981]. Firstly, the low frequency wave energy of a large amplitude wave may be confined to the small volume of the soliton yielding intense interaction. Secondly, the soliton may collapse thereby enhancing this effect due to increase of the wave energy density [Treumann, 1982]. Thirdly, in restricting the acceleration to the soliton volume, the soliton starts effectively acting as an antenna; hence emission is not determined by the high frequency waves occupied by the soliton but by the volume filled with high frequency waves which the soliton traverses during the interaction time and its motion across the plasma. Finally, emission by solitons should yield narrow band radiation because of the finite spatial extension of the solitons along field lines. Moreover, several different frequencies and modes are emitted in this process due to the participation of the trapped mode, the envelope mode, and Doppler shift due to the soliton motion, while the modes arise due to longitudinal as well as transverse components of the oscillating currents generated in the acceleration process. The former lead to O–mode, the latter to X–mode emission. In addition, one expects highly variable radiation from this process because the solitons tend to evolve in time and to collapse. The requirement of participation of high frequency waves renders the process highly sensitive on the excitation mechanisms of both waves participating in the interaction.

In the following we restrict to only the simplest case: Nearly perpendicular propagation of the lower hybrid waves, emission in the X–mode only, and neglect of all weakly excited modes. We also assume a stationary generation mechanism of as well lower as upper hybrid waves arguing that both waves can exist coincidentally in the same region of space and at the same time when they are generated by electron beams. Conditions where this can be the case have previously been given by Goldstein et al. [1983] who also advocated wave coupling as a possible radiation mechanism.

Figure 1 is a sketch of the model. It shows the idealized lower hybrid soliton elongated along the magnetic field lines and propagating with velocity $\mathbf{u}_\perp$ perpendicular to the magnetic field. Actually, the soliton can as well have a finite velocity $u_\parallel$ along the magnetic field. The soliton is embedded into the path of upper hybrid waves shown as sinusoidal wave trains. Electrons accelerated in both fields emit long wavelength radiation in the X–mode.

3 Radiated power

The power emitted by the accelerated electrons is given by

$$\frac{\mathrm{d}\mathcal{P}}{\mathrm{d}\Omega} = \frac{1}{\mu_0} k_0^2 R^2 c |A(\omega)|^2 \left(1 - \left[\frac{\mathbf{n}\cdot\mathbf{A}}{|A|^2}\right]^2\right), \tag{3.1}$$

where $\mathcal{P}, \Omega, A(\omega)$ are the respective power, solid angle, and vector potential of the radiation, ω, k_0 are frequency and wave number of the radiation, $\mathbf{n}$ is the direction in which radiation is emitted, and R is the distance from the source to the observer.

The vector potential if taken in the far field approximation can be represented by

$$\mathbf{A}(t) = \frac{\mu_0 \exp(ik_0 R)}{4\pi R} \int \mathbf{j}^{nl} \mathrm{e}^{-k_0 \mathbf{n}\cdot\mathbf{r}'} \mathrm{d}^3 r' \,. \tag{3.2}$$

It contains the nonlinear electron current density $\mathbf{j}^{nl}$ which using electron momenta is given by the ensemble averaged product of the fluctuation in density and field. It is convenient to use its Fourier representation

$$\mathbf{j}^{nl}(\omega, \mathbf{k}) = -e \int \delta n(\omega', \mathbf{k}') \mathbf{v}(\omega - \omega', \mathbf{k} - \mathbf{k}') \mathrm{d}\omega' \mathrm{d}^3 k' \,. \tag{3.3}$$

We have introduced the variations in density δn and velocity $\mathbf{v}$ of the electrons assuming that there is no zero order streaming of the electron distribution. The latter would only complicate the equations but not contribute to radiation.

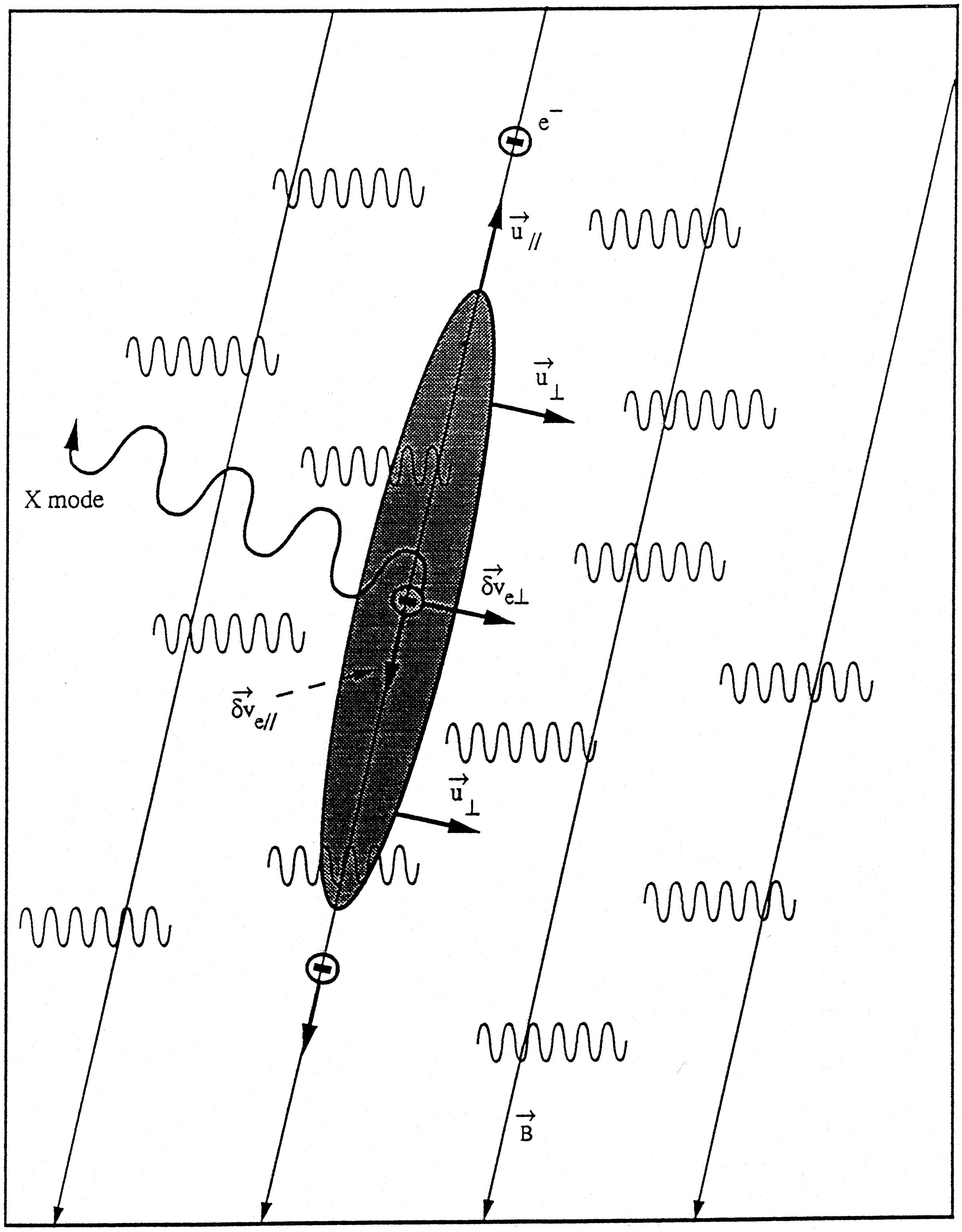

Figure 1: Sketch of the soliton radiation mechanism. The lower hybrid soliton is shadowed. In its initial stage it is elongated along the magnetic field lines and propagates on the path of high frequency upper hybrid waves at velocity **u**. *Both types of waves are oblique to the magnetic field. Electrons entering the region of the soliton are accelerated in the combined electric field and radiate an X–mode wave perpendicular to the magnetic field.*

4 Oscillating nonlinear current

The form of the nonlinear oscillating source current which produces the emitted radiation depends heavily on the underlying model used in calculating the electron momenta. In a kinetic approach one has to solve the Vlasov equation for electrons moving in the combined fields of the lower hybrid soliton and the upper hybrid wave. Since the electrons gain or lose energy in this process, the calculation would have to be performed selfconsistently taking into account a number of processes as wave excitation, soliton formation and damping by the interaction, emission and reabsorption of the radiation.

A much simpler approach is to pass straight to an approximate hydrodynamic collisionless model, where the electrons are considered a fluid and the ions a neutralizing otherwise stationary background. Actually the ions contribute to the generation of the lower hybrid waves and solitons [Gekelman and Stenzel, 1975; Musher et al., 1978], but this contribution can be considered separately and independent of the radiation process as long as the loss of energy of the lower hybrid soliton in the interaction with the radiating electrons can be neglected. At the high radiative frequencies this is a reasonable approach since in first approximation the electrons move adiabatically in the lower hybrid field, and it is basically the nonlinear phase mixing which mixes the frequency and hence part of the energy of the low frequency wave field into the radiation to enable it to overcome the X–mode cut–off.

In the hydrodynamic model we linearize the electron equations of continuity and motion, define $\Delta(\omega) = \Omega_e^2 - \omega^2$ and make use of the cold plasma mobility tensor $\mathcal{M}$ to obtain

$$\begin{aligned}\mathbf{j}^{nl}(\omega, \mathbf{k}) = -en_0 \quad & \int \tfrac{\mathrm{d}\omega'}{\omega'} \mathrm{d}^3k'\mathbf{k}' \cdot [\mathcal{M}(\omega') \cdot \mathbf{E}(\omega', \mathbf{k}')] \\ & \times \mathcal{M}(\omega - \omega') \cdot \mathbf{E}(\omega - \omega', \mathbf{k} - \mathbf{k}') \,. \end{aligned} \tag{4.1}$$

As explained above we neglect any thermal effects on the electrons. These complicate the mobility tensor introducing an additional dependence on k' and consequently enhance the number of possible emission bands. Effects like this could more appropriately be taken into account in a kinetic treatment.

For given electric wave fields, Equation (4.1) solves the problem. We assume an external magnetic field $\mathbf{B} = B\mathbf{z}$ along the z–axis and the wave vectors of the two wave modes participating in the acceleration to be in the xz–plane. The total electric field experienced by the electron is the sum of the high frequency field of the upper hybrid wave ($\mathbf{E}_h$) and the low frequency field of the lower hybrid soliton ($\mathbf{E}_l$). Since the upper hybrid wave is plane electrostatic with potential $\phi_h(\mathbf{r}, \mathbf{t}) = \phi_u \exp i(-\omega_{\mathbf{k}h}t + \mathbf{k}_h \cdot \mathbf{r}) + c.c.$ one has

$$\mathbf{E}_h(\mathbf{k}, \omega) = -i\mathbf{k}\phi_u \delta(\omega - \omega_h)\delta(\mathbf{k} - \mathbf{k}_{\perp h}) \,, \tag{4.2}$$

where $\omega_h = \omega_{\mathbf{k}h} \cdot \mathbf{u}$ and we have transformed to the moving frame of the soliton.

The soliton field can be written as $\phi_l(\mathbf{r}, t) = \varphi(\mathbf{r}, t) \exp i(-\omega_{\mathbf{k}l}t + \mathbf{k}_l \cdot \mathbf{r}) + c.c.$, where the envelope φ_l is a function which varies slowly over the spatial and temporal scales of the lower hybrid wave and is given by

$$\varphi_l(\mathbf{r}, t) = \varphi_l \mathrm{sech}\left(\frac{x}{L_\perp} - \frac{z}{L_\parallel}\right) \exp i(-\omega_s t + \mathbf{k}_s \cdot \mathbf{r}) \,. \tag{4.3}$$

The subscript s indicates the small frequency and momentum shifts of the soliton and $L_{\|,\perp}$ are the parallel and perpendicular widths of the soliton. We have $|k_s| \ll |k_l|, \omega_s \ll \omega_l$ and the contributions of the shifts to the bands are small and will be neglected. One observes, however, that they contribute to small splittings of the emission bands. Using (4.2) and (4.3), one finds an expression for the nonlinear oscillating current. In general this current has three components and several terms contributing to each of the components. Since the perpendicular wave amplitudes of the electrostatic waves involved are much larger than the parallel components, the dominant contribution comes from the perpendicular components of $\mathbf{E}$ and $\mathbf{k}$. One finds that the dominant term oscillates at frequency $\omega \approx \omega_l + \omega_h$ yielding a radiation vector potential of the form

$$
\begin{aligned}
A_x &= D\mu_0\pi^{\frac{3}{2}}\frac{\exp(ik_0R)}{R}L_n(k_{hx}\omega_l + k_{lx}\omega_h) \\
&\times\mathrm{sech}\left(\frac{\pi}{2}L_\perp\alpha\right)\exp\left(-\frac{k_0^2L_n^2}{4}\sin^2\theta\sin^2\psi\right) \\
&\times\delta(\omega-\omega_l-\omega_h)\delta(\alpha L_\perp + \beta L_\|)\,, \qquad (4.4)\\
A_y &= \frac{i\Omega_e(k_{lx}+k_{hx})}{k_{lx}\omega_h + k_{hx}\omega_l}A_x(\omega)\,, \\
A_z &\simeq 0\,.
\end{aligned}
$$

Here

$$
\begin{aligned}
\alpha &= k_{lx} + k_{hx} - k_0\sin\theta\cos\psi\,, \\
\beta &= k_{lz} + k_{hz} - k_0\cos\theta\,, \qquad (4.5)
\end{aligned}
$$

and

$$
D = \frac{e^3n_0}{m_e^2}\frac{L_\|L_\perp k_{lx}k_{hx}\varphi_l\varphi_h}{(\Omega_e^2-\omega_l^2)(\omega_h^2-\Omega_e^2)}\,. \qquad (4.6)
$$

The main contribution to the vector potential comes hence from the x–component. Further, the Delta–function of ω determines the frequency of the radiation which is found to be the predicted fundamental radiation frequency $\omega \approx \omega_h + \omega_l$. This expression corresponds to the resonance condition between the waves. It is most important to ask under which conditions this expression can be satisfied and leads to escaping radiation.

The second Delta–function depends on the wave vector components yielding the condition

$$
L_\perp/L_\| \sim -\beta/\alpha\,. \qquad (4.7)
$$

Because for lower hybrid solitons $L_\| \gg L_\perp$, this condition implies $\alpha \gg |\beta|$. At the same time the sech–dependence of the radiation vector potential suggests a mismatch in the perpendicular wave number of the order of $\Delta k_\perp \sim 2/\pi L_\perp$. The wave number condition is hence not a sharp resonance as required in wave coupling theories [cf. e.g. Goldstein et al., 1983]. Neglecting k_0 and using the soliton parameters (see Section 5), the parallel wave number mismatch becomes

$$
|\Delta k_\||\lambda_i \sim \frac{2}{\pi\sqrt{3}}\left(\frac{m_e}{m_i}\right)^{\frac{1}{2}}\,. \qquad (4.8)
$$

There is hence a wide range of parallel and perpendicular wave numbers which can satisfy the radiation conditions, in remarkable contrast to the wave coupling theories. On the other hand, since the radiation wave length is long, the wave vectors of the upper and lower hybrid wave fields must be of opposite sign but can differ by the above amounts in parallel and perpendicular direction.

The emission condition requires that the emitted frequency exceeds the low frequency X–mode cut–off $\omega_\times$ where

$$\omega_\times \approx \Omega_e + (\omega_e^2/\Omega_e)\,. \tag{4.9}$$

The dispersion relation of the upper hybrid waves [Porkolab, 1984] reads

$$\omega_h \approx \omega_u[1 - (k_\parallel/k)^2\omega_e^2\Omega_e^2/\omega_u^4]^{\frac{1}{2}}\,. \tag{4.10}$$

This frequency is close to the upper hybrid frequency ω_u. At the same time the dispersion relation of the lower hybrid wave, neglecting the contribution of the soliton frequency and wave number shift, is given by

$$\omega_l \approx \omega_{lh}(1 + \frac{m_i}{m_e}\frac{k_\parallel^2}{k^2} + 3k^2\lambda_i^2)^{\frac{1}{2}}\,, \tag{4.11}$$

with ω_{lh} the lower hybrid resonance frequency.

In the nonlinear state, when solitons have been formed, one has for the lower hybrid wave that $k_l\lambda_i \sim 1$. The maximum growing upper hybrid waves, on the other hand, have $k_h\rho_e \sim 1$ where the warm electron gyroradius ρ_e has been introduced. Approximate equality of these two expressions yields then that $\omega_e/\Omega_e \sim v_e/v_{eh}$, with v_{eh} the hot electron temperature. In the case of exact resonance as in wave coupling theory this expression would have to be exactly equal thereby restricting the interaction to a narrow range where the ratio of electron plasma to cyclotron frequency becomes identical to the ratio of background to hot electron temperatures. Emission would then take place only in a narrow region where these conditions are satisfied. As obvious the condition is valid only in a dilute plasma. In the case of soliton interaction the condition is largely relaxed due to the finite wave number range which is allowed to participate in the radiation process. This fact reflects the catalytic action of the soliton which simply serves as an antenna for transforming high frequency wave energy into radiation without contributing essentially to the resonance.

Checking the escaping condition one must take into account that the electrostatic waves are slightly oblique with respect to the magnetic field. It is sufficient to require that, for the conditions when for instance as in the polar magnetosphere the ratio of plasma to cyclotron frequency is of the order of $\omega_e/\Omega_e \approx 0.15$ [Pottelette et al., 1990], the wave vector $\mathbf{k}_l$ of the lower hybrid waves deviates by about 3° from the perpendicular to the magnetic field to guarantee that the escaping radiation frequency $\omega_0 > \omega_\times$ exceeds the X–mode cut–off. Goldstein et al. [1983] have investigated the similar condition for Jupiter wave coupling emission and found a similar result. Using the wave number mismatch conditions even this relation can be relaxed because the finite $\Delta k_\parallel$ inserted into the frequency matching law $\omega_0 = \omega_l + \omega_h$ contributes to raising the radiated frequency above the cut–off limit. As a consequence, the emitted frequency in this process will be well above the X–mode cut–off even for nearly transverse propagation of the lower hybrid waves involved in the solitons. Such oblique lower hybrid waves are of course closely related to resonant whistlers and will also have an electromagnetic component which has been neglected here.

5 Soliton dynamics

The dynamics of lower hybrid solitons is a difficult chapter. It is not the place here to go into it in depth. The solitons which are needed for making the mechanism work, have been assumed low frequency because the observations have shown that the emission is very close to the fundamental frequency, in our case of low plasma to cyclotron frequency ratios to the electron cyclotron frequency itself. The previous discusssion dismantled on the other hand that low frequency solitons are well suited to generate radiation near but sufficiently high above the fundamental that the waves can escape.

In brief, lower hybrid solitons can form when for instance lower hybrid waves are generated by fast electrons flowing along the magnetic field lines [e.g. Goldstein et al., 1983] or when anisotropic or loss cone distributions in the warm electron population generate hiss in the whistler mode [e.g. Lotko and Maggs, 1979; Chang and Coppi, 1981] and part of this whistler spectrum transforms near the resonance cone into lower hybrid waves. Especially, in the auroral magnetosphere Maggs and Lotko [1981] have shown that hiss waves readily turn their wave vector perpendicular to the magnetic field when propagating down to the ionosphere thereby becoming resonant and electrostatic and making the transition to lower hybrid waves. Convective losses are important for these waves, but when the beam to thermal density ratio becomes $n_b/n_0 \gtrsim 10^{-3}$ nonlinear amplitudes are reached long before the waves can be convected out of the unstable region. These waves then can evolve into solitons [Sturman, 1976] when the waves are scattered at electrons, lose momentum and condensate at long wave lengths so that

$$3k^2\lambda_i^2 < (m_i/m_e)k_\parallel^2/k^2 \,, \tag{5.1}$$

in which case the character of the dispersion of the lower hybrid waves changes and solitons develop. The evolution of the solitons has been discussed in Pottelette et al. [1992]. The wave energies required for soliton formation are small, of the order of $W_l/nT_e \lesssim m_e/m_i$. This condition holds in the regime where the cyclotron frequency is much larger than the electron plasma frequency. Hence this regime favours soliton formation. Measured wave levels are of the order of $W_l/nT_e \approx 0.003$ only about a factor 5 higher than the above limit. It has also been shown that quasilinear saturation of the beam exciting the waves takes about one order of magnitude longer than modulation and soliton formation for beams of same temperature as the background. Since quasilinear time is proportional to the ratio T_{eb}/T_e hot beams which are expected to exist under planetary auroral conditions never saturate quasilinearly in lower hybrid resonance.

Solitons on the other hand will not exist forever but will either decay due to interaction with fast electrons and due to radio wave emission or will collapse. There is one curious fact about lower hybrid solitons. They are by nature two–dimensional and in addition and in contrast to Langmuir solitons do not conserve the full plasmon number trapped in the solitons during collapse. As a consequence, plasmons leak out from the solitons during collapse and may form other daughter solitons. In this way a whole family of solitons may be formed and the plasma may become nonlinearly structured assuming a turbulent state with large fluctuations in density and wave field even if only one soliton has developed initially.

The dimensions of the soliton parallel and perpendicular to the magnetic field are related via $L_\parallel/L_\perp \approx (3m_i/m_e)^{1/2}k\lambda_i \gg 1$. On the other hand, the relation between the soliton amplitude A and the transverse dimension becomes

$$AL_\perp \approx 6\sqrt{2}\lambda_i[T_i(T_e+T_i)]^{\frac{1}{2}}/e\,, \tag{5.2}$$

or using $k\lambda_i \sim 1$ and $T_e = T_i$ approximately $|E_l|L_\perp \approx 12T_e/e$. Using a measured maximum value of 100 mV/m [Pottelette et al., 1990] and $T_i \approx 10$ eV, one finds a characteristic transverse width of the soliton of $L_\perp \sim 1$ km. Note that this size does excellently agree with the tuning condition of an antenna for radiating km–wavelengths radio waves.

During collapse the two dimensions of the soliton evolve differently according to

$$L_\perp \propto \lambda_i(me/m_i)^{\frac{1}{2}}(\Omega_e/\omega_e)(nT_e/W_l)^{\frac{1}{2}}\,, \tag{5.3}$$

and

$$L_\parallel \propto (\sqrt{3}\lambda_i)^{-1}(m_i/m_e)^{\frac{1}{2}}L_\perp^2\,. \tag{5.4}$$

From these equations one finds that $L_\perp \propto (1-t/t_c)^{1/2}$ and $L_\parallel \propto (1-t/t_c)$ where t_c is the collapse time. According to this the parallel dimension collapses faster than the perpendicular. The soliton in its final stage will then become a pancake, and radiation will cease quickly because of the shrinkage of the radiating volume.

6 Estimates

As long as the soliton exists radiation can be emitted from its region. Using the above expressions one finds that the bandwidth of the emission at the fundamental $\omega \approx \omega_h + \omega_l$ is of the order of

$$\Delta f \sim \frac{3L_\parallel}{H} f_{ce}\,, \tag{6.1}$$

where we used a dipole magnetic field and H is the planetocentric altitude. For the Earth we find ($H \approx 1.9R_E, f_{ce} \approx 250$ kHz) a value of about $\Delta f \approx 2$ kHz corresponding to the observed bandwidth of the elements of radiation in the AKR [Gurnett et al., 1981; Benson and Calvert, 1979].

Under the same conditions, assuming a soliton volume $V_s = L_\perp L_\parallel L_n$ (L_n is the density gradient scale assumed much larger than the dimension of the soliton), of the order or some ten km^3, an upper hybrid wave field amplitude $E_h \approx 3$ mV/m, we find a power of $\mathcal{P} \approx 13$ W per soliton. Because of the uncertainties involved, the power may hence range between 1–100 W per soliton. This corresponds to fluxes of $10^{-11} - 10^{-9}$ Wm^{-2}Hz^{-1} in approximate agreement with observation. To account for the total observed AKR power one needs about 10^6 solitons in the source volume. In the more favourable case these solitons fill a volume of 10% of the available AKR source volume. This condition might not always be satisfied implying that only part and not all of the emitted AKR power is generated by soliton radiation.

7 Discussion

The mechanism proposed here is based on the nonlinear evolution of lower hybrid waves in the dilute polar plasma of magnetized planets. In application to the AKR of the Earth it yields sufficiently high radiated powers to account for part of the radiated emission. At the same time it yields rather narrow elementary emission bands from one single soliton. The broad planetary radio emission must then be caused by a large number of elementary solitons moving on the background of a distribution of upper hybrid waves. Since the solitons themselves add only little energy to the radiation, the energy is taken from the high frequency electrostatic upper hybrid waves. Hence the mechanism heavily depends on the simultaneous presence of both types of waves in the source volume. Changes in these conditions will be immediately reflected in changes in the emission, its bandedness and duration. Hence, this mechanism because of its high sensitivity to such variations will be highly variable in time and space. It has however the advantage that the radiation conditions are not as restrictive as for wave–wave coupling where no resonance broadening is allowed for. The soliton mechanism allows for participation of a broad k–spectrum of waves in the radiation process.

The relatively high energy a single soliton emits is due to its motion across the background of high frequency waves. Since interaction takes place only inside the soliton, the moving soliton works like a vacuum cleaner eating up all the available high frequency waves on its pass through the source region (Figure 2). As one observes, one therefore does not require very high energies to be present in the single upper hybrid waves to obtain moderate radiation intensities. Moreover, the mechanism does not provide problems in exceeding the low frequency cut–off of the X–mode.

In addition to X–mode radiation there are also terms contributing to O–mode emission. The intensities of these emissions are, however, weaker than X–mode intensity. Nevertheless it is important to note that O–mode has been excluded by us from the discussion only because of its weakness. Its presence is not included in general. Further investigation of the mechanism in this direction is of interest inasfar as O–mode radiation has been reported in the past.

The emission does actually consist of many different bands which are separated by small contributions to the emitted frequency provided by the different frequency and wave number shifts involved in the soliton formation and the interaction process. All these contributions contribute to fine structure. However, only part of this fine structure could be resolved by present observational methods.

In conclusion, we have presented a simple alternative mechanism for generation of electromagnetic radiation from a dilute plasma. This mechanism requires similar outer conditions as the well–known maser radiation mechanism. One can therefore expect that it

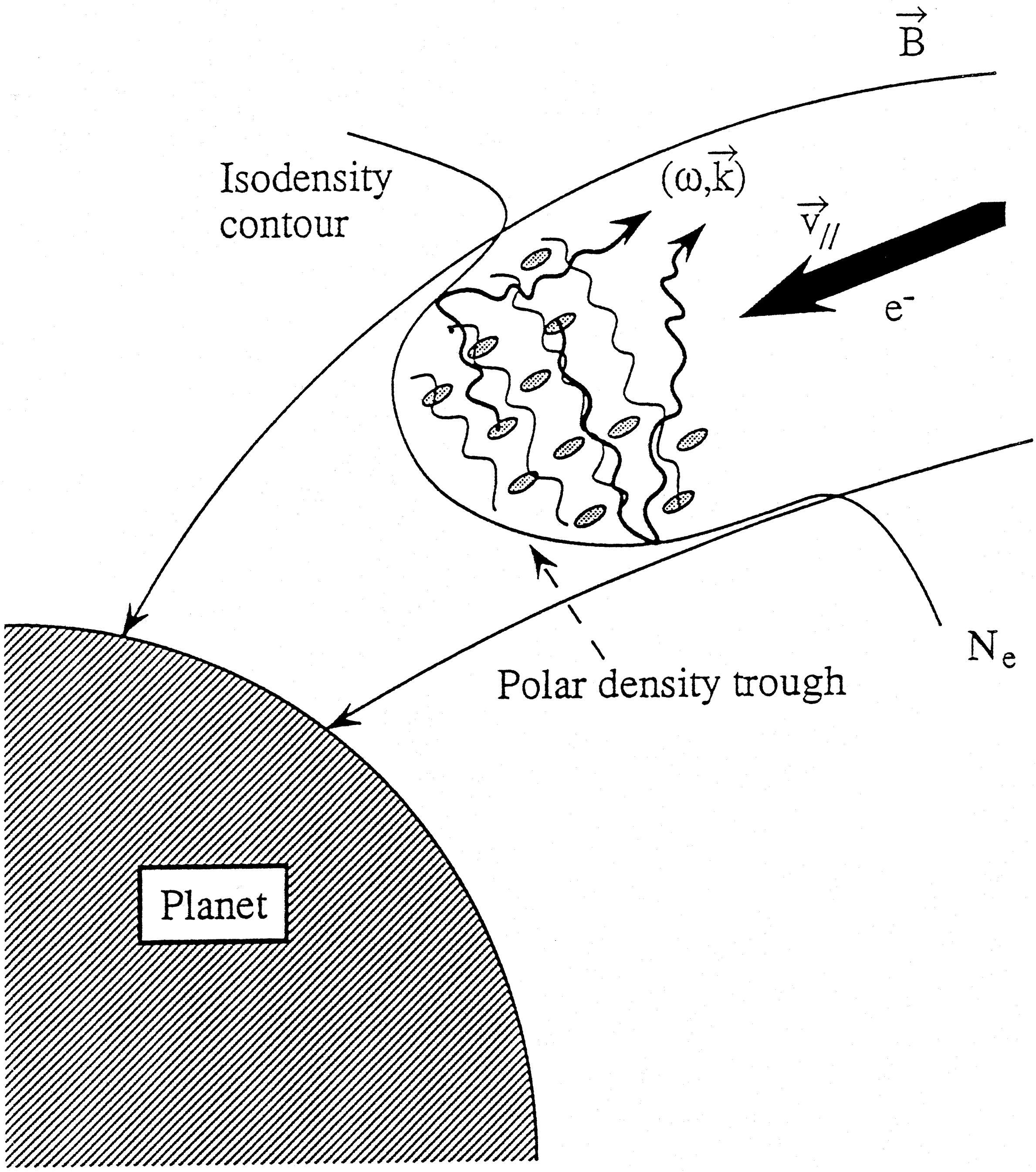

Figure 2: Sketch of the source region for planetary soliton antenna radiation. The isodensity contour indicates presence of a density depletion region filled with upper hybrid waves and lower hybrid solitons. Emission from the solitons is shown to take place. It may become reflected at the trough boundary before leaving the source region. The heavy arrow designs fast electrons flowing down the field lines. The sketch also suggests that reabsorption of radiation is possible when radiation on its way out hits some other soliton.

may work in close connection with the latter contributing to the emission. It does however neither replace nor exclude maser radiation wherever the condition are in favour of the latter simply because maser emission is a direct linear process which yields high power. It however compliments the maser mechanism in some way. One can also anticipate the existence of a relation or coupling between both processes. Especially the escape of maser radiation may provide some problems. In this case the maser emission may when interacting with the boundaries of the dense region provide the required plasma waves which then start acting as antennas. Another possibility is that reabsorption of the emission of one or the other kind in the plasma stimulates emission in the other mechanism or mode. Further investigation should analyse the relation between these two basic radiation mechanisms. It should also investigate in greater detail the fine structure in frequency and time provided by the soliton antenna emission mechanism.

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